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## Domain Wall World(s)

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### ABSTRACT

Gravitational properties of domain walls in fundamental theory and their implications for the trapping of gravity are reviewed. In particular, the difficulties to embed gravity trapping configurations within gauged supergravity is reviewed and the status of the domain walls obtained via the breathing mode of sphere reduced Type IIB supergravity is presented.

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<sup>1</sup>Based on the talk given at Strings'2000.

Over the past few years domain walls have played an important role both from the point of view of the AdS/CFT correspondence, i.e. by shedding light on the renormalization group flow and bound state spectra of strongly coupled field theories, as well as from the point of view of the brane world scenarios, such as Randall-Sundrum scenario [1] for localization of gravity on domain walls in five-dimensions. In this contribution, however, we shall focus on gravitational properties of domain walls, their implications for the brane world scenarios and their implementation in the fundamental theory.

This presentation is based on the work with K. Behrndt [2, 3], along with the parallel work by Kallosh and Linde and Schmakova [4, 5], leading to the “no-go” theorems for the implementation of the Randall-Sundrum scenario, when one employs massless modes of gauged supergravity. (For subsequent developments see, e.g., [6, 7, 8].) The focal part, however, will be based on the work with H. Lü and C. Pope [9, 10] where the massive-breathing mode of sphere reduced gauged supergravity was proposed as the candidate field for the gravity trapping domain walls. For subsequent related works see [11] and in particular [12] as well as the contributions to these proceedings by S. de Alwis, J. Liu and K. Stelle. In addition, there the origin of the Randall-Sundrum solution [13] based on  $S^1/Z_2$  topology and two branes sources was addressed in detail.

The outline of this contribution is the following: (i) We shall first review the properties of supergravity domain walls, identifying the flat (Minkowski) walls with BPS saturated ones and the bent (deSitter and anti-deSitter) ones with those whose energy density is either larger or smaller than that of the BPS saturated ones. (This part of the review is primarily based on a much earlier work on domain walls in four-dimensional N=1 supergravity, initiated in [14], its generalizations to non-BPS solutions were given in [15] and reviewed in [16]; for a recent work on a generalization of the analysis to D-dimensions see [17] and references therein.) (ii) We shall then turn to the discussion of the properties of domain walls in N=2 D=5 gauged supergravity (with vector multiplets), leading to the “no-go” theorem for the implementation of the Randall-Sundrum scenario. (iii) As the last topic we turn to the status of the breathing mode domain walls, where the breathing mode parameterizes the volume of the compactified Einstein space (with the sphere as the most symmetric example). While these walls may provide the candidate gravity trapping solutions, we also mention difficulties with the interpreting of the delta function sources for infinitely thin domain wall within Type IIB string theory.

# 1 Domain walls in N=1 supergravity

Since the scalar potential in supergravity theories is of the restricted form it is natural to expect that the nature of domain walls is of special form as well. The bosonic Lagrangian of D=4 N=1 supergravity theory takes the form:

$$\mathcal{L} = g_{A\bar{B}} \partial_\mu \Phi^A \partial^\mu \Phi^{\bar{B}} - V - \frac{R}{8\pi M_{pl}^2}, \quad (1)$$

where the scalar potential (for the gauge neutral fields) is of the form:

$$V = g^{A\bar{B}} \partial_A \hat{W} \partial_{\bar{B}} \hat{W} - \frac{3}{M_{pl}^2} \hat{W}^2. \quad (2)$$

Here the Kähler invariant quantity  $\hat{W}$  takes the form:  $\hat{W} = \zeta |W| e^{\frac{K}{2M_{pl}^2}}$ , where  $\zeta = \pm 1$  and  $\zeta$  can change the sign only when  $W$  goes through the zero value. Here  $W$  and  $K$  are the superpotential and Kähler potential, respectively. Supersymmetric extrema are at  $\partial_A \hat{W} = 0$ , and correspond to either a zero cosmological constant, i.e.  $\Lambda \equiv V_{ext} = 0$  when  $\hat{W}_{ext} = 0$ , or a negative cosmological constant, i.e.  $\Lambda \equiv V_{ext} < 0$ , when  $\hat{W}_{ext} \neq 0$ .

The surprising result [14] in the study of domain walls in supergravity theory is that the static (flat) domain wall solutions between supersymmetric extrema do exist irrespective of the fact that the actual value of the cosmological constants of the two isolated extrema may be different. Such solutions turn out to be BPS ones: they satisfy the Killing spinor equations:  $\delta\psi^\alpha = 0$ , and  $\delta\psi_\mu^\alpha = 0$  for the supersymmetric variation of the scalar-field superpartner and gravitino, respectively. These equations are solved with the static Ansatz for the scalar field  $\Phi$  (only one chosen for simplicity) and the following conformally flat metric Ansatz<sup>2</sup>:

$$ds^2 = A(z)(-dt^2 + dx^2 + dy^2 + dz^2) \equiv \mathcal{A}(\tilde{z})(-dt^2 + dx^2 + dy^2) + d\tilde{z}^2, \quad (3)$$

leading to the following coupled first order equations of motion [14]:

$$\begin{aligned} \partial_z \Phi^A &= 2g^{A\bar{B}} \partial_{\bar{B}} \hat{W}, \\ M_{pl}^2 \partial_z \log \mathcal{A} &= -2\hat{W}, \end{aligned} \quad (4)$$

and the energy density of the wall takes the following form:

$$\sigma_{BPS} = 2(\hat{W}_+ - \hat{W}_-) = 2M_{pl}(\zeta_+ \sqrt{-\frac{\Lambda_+}{3}} - \zeta_- \sqrt{-\frac{\Lambda_-}{3}}). \quad (5)$$

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<sup>2</sup>We shall describe the domain wall metric interchangeably in terms of the conformal factor  $A$  or the warp factor  $\mathcal{A}$ .

The asymptotic behavior of the solution is the following:

$$\begin{aligned} z \rightarrow \pm\infty : \quad \Phi &\rightarrow \Phi_{\pm} + e^{2\partial_{\Phi}^2 \hat{W}_{\pm} z}; \quad \{\partial_{\Phi}^2 \hat{W}_{-} > 0, \partial_{\Phi}^2 \hat{W}_{+} < 0\}, \\ \mathcal{A} &\rightarrow e^{-\frac{2\hat{W}_{\pm} z}{M_{pl}^2}}. \end{aligned} \quad (6)$$

Note that  $\text{sign}(\hat{W}_{\pm})$  determines the asymptotic behavior of the metric. It turns out that the necessary condition for the exponential fall-off of the metric warp factor  $\mathcal{A}$  is that at the supersymmetric extrema the potential  $V(\Phi_{\pm})$  has the minimum.

A typical  $Z_2$  symmetric example corresponds to the following choice:

$$W = \sqrt{\lambda}\Phi\left(\frac{\Phi^2}{2} - \eta^2\right), \quad K = \Phi\Phi^*, \quad (7)$$

where the kink solution (4) interpolates between two supersymmetric anti-deSitter minima with the same cosmological constant ( $\hat{W}_{+} = -\hat{W}_{-} \equiv \hat{W}_0 = M_{pl}\sqrt{-\frac{\Lambda}{3}}$ ). The metric conformal factor  $A(z)$  (3) falls-off on either side as  $M_{pl}^4(\hat{W}_0 z)^{-2}$ , and the energy density of the wall is  $\sigma_{BPS} = 2 \times 2\hat{W}_0 = 4M_{pl}\sqrt{-\frac{\Lambda}{3}}$ .

The thin wall limit of the  $Z_2$  symmetric BPS solution is achieved by taking the limit  $\lambda \rightarrow \infty$ ,  $\eta \rightarrow 0$ , while  $\lambda\eta^3$  remains fixed. In this particular case the superpotential  $\hat{W}$  approaches the step function, while the potential and the scalar kinetic energy term approach the delta function. This leads to the thin wall effective (supersymmetric) Lagrangian with the delta function source that precisely reproduces the thin-wall flat domain-wall solution, discussed by Randall-Sundrum in D=5. (For further discussions for a supersymmetric implementation of the effective Lagrangian with the delta function sources see [18, 19, 20, 21, 22].)

In the case of broken supersymmetry, the walls become bent. In the thin wall analysis of walls that can assume the  $Z_2$  symmetric limit, the space-time internal to the wall is either deSitter (when  $\sigma > \sigma_{BPS}$ ) or anti-deSitter (when  $\sigma < \sigma_{BPS}$ ). The conformal factor (3) on the other hand takes the form [15]:

$$\begin{aligned} A(z) &= M_{pl}^2 \beta^2 \left[ \sqrt{-\frac{\Lambda_{\pm}}{3}} \sinh(\beta(z + z_{\pm})) \right]^{-2}, \quad \text{for } \sigma > \sigma_{BPS}, \\ A(z) &= M_{pl}^2 \beta^2 \left[ \sqrt{-\frac{\Lambda_{\pm}}{3}} \cosh(\beta(z + z_{\pm})) \right]^{-2}, \quad \text{for } \sigma < \sigma_{BPS}. \end{aligned} \quad (8)$$

where  $\beta^2 \sim |\Lambda_{wall}|$  specifies the cosmological constant on the bent wall for respective anti-deSitter and deSitter space-times on the wall. The constants  $z_{\pm}$  are appropriately chosen so that at  $z = 0$  the conformal factor is  $A(z)$  is normalized to 1. In the former case the conformal factor falls-off even faster than in the BPS case ( $|z| \rightarrow \infty$  corresponds to the cosmological horizons), while for the latter case the conformal factor, while first decreasing

it turns around at intermediate distances ( $|z| \rightarrow \infty$  is the space-time boundary). This latter case allows for the possibility of quasi-localized gravity on the wall [23].

While we have briefly reviewed the BPS and non-BPS thin walls for supergravity theory in four dimensions, it was shown in [17] that the global and local space-time structure of co-dimension one objects in D dimensions is universal, only that the role of the cosmological constant factor  $\sqrt{-\frac{\Lambda_{\pm}}{3}}$  is replaced by  $\sqrt{-\frac{(D-2)\Lambda_{\pm}}{2(D-1)}}$ . In particular, for domain walls in D=5, this factor becomes  $\sqrt{-\frac{3\Lambda_{\pm}}{8}}$ .

## 2 Domain Walls in D=5 Gauged Supergravity

A natural question to be asked is to identify the origin of gravity trapping solutions in fundamental (M-)theory, as initiated in [24]. In particular, can such domain walls arise in an effective five-dimensional theory, that can be obtained as a compactification of, e.g., M-/string theory on Einstein-Sasaki spaces? Since such compactifications are expected to produce an effective gauged supergravity theory, we now turn to the study of domain walls in five-dimensional N=2 gauged supergravity. For the sake of simplicity, we shall focus on D=5 N=2 gauged supergravity with Abelian  $(U(1)_R)$ -gauging and with the vector supermultiplets, only [25].

The bosonic sector is of the constrained form, with the vector supermultiplets  $X^I$  (real, neutral fields) subject to the following condition:

$$F = \sum_{IKJ} C_{IKJ} X^I X^J X^K = 1, \quad (9)$$

which can be solved for the physical scalar fields  $\Phi^A$ . The potential of  $U(1)$  gauged supergravity is also of the constrained form:

$$W = \sum_I h_I X^I. \quad (10)$$

The bosonic part of the Lagrangian nevertheless takes an analogous form as in the case of D=4 N=1 supergravity (We have now set  $M_{pl}=1$ .):

$$\mathcal{L} = g_{AB} \partial_\mu \Phi^A \partial^\mu \Phi^B - V + R, \quad (11)$$

where the metric and the potential are of the form:

$$g_{AB} = \frac{1}{2} (\partial_I \partial_J F) \partial_A X^I \partial_B X^J; \quad V = g^{AB} \partial_A W \partial_B W - \frac{4}{3} W^2. \quad (12)$$

Again, the Killing spinor equations, corresponding to the supersymmetric domain wall solutions, reduce to the equations of the analogous type:

$$\begin{aligned}\partial_z \Phi^A &= \mp 3g^{AB} \partial_B W, \\ \partial_z \log \mathcal{A} &= \pm 2W.\end{aligned}\tag{13}$$

However, due to the fact that  $W$  and  $g_{AB}$  are of the constrained form, they satisfy a relationship [25], that takes the following form at the supersymmetric extrema:

$$\partial_A \partial_B W_{ext} = \frac{2}{3} [g_{AB} W]_{ext}.\tag{14}$$

Expanding the Killing spinor equations around the supersymmetric minima  $\partial_\Phi W|_\pm = 0$ , as  $\Phi^A = \Phi_\pm^A + \delta\Phi^A$  and using the relationship (14), yields the following asymptotic form of these equations:

$$\begin{aligned}\partial_z(\delta\Phi^A) &= \mp W_\pm \delta\Phi^A, \\ \partial_z(\log \mathcal{A}) &= \pm W_\pm.\end{aligned}\tag{15}$$

The condition for the asymptotic kink solution requires  $\text{sign}(W_+) = -\text{sign}(W_-)$ , and then the conspiracy of signs in the above equations (15) necessarily requires that the kink solution has the metric factor that *grows* exponentially on either side of the wall. Thus, these solutions are not relevant for the trapping of gravity; those are typical domain walls relevant for AdS/CFT correspondence. In addition, the constraint (14) implies that supersymmetric extrema with  $W_{ext} > 0$  are necessarily the minima of the superpotential (for  $g_{AB}$ -positive definite), while supersymmetric extrema with  $W_{ext} < 0$  are maxima. Therefore, the kink necessarily has to cross the singular region in the superpotential manifold and thus the solution is generically singular.

Further studies of more general solutions reveal that neither inclusion of non-Abelian tensor multiplets [6, 5], nor inclusion of hypermultiplets [7, 8] allow for non-singular domain walls that would have a fall-off metric conformal factor on both sides of the wall. (For a related no-go theorem see [26].)

### 3 Breathing mode and gravity trapping domain walls

In this section we shall review a framework within gauged supergravity theories that has a chance of implementing a variant of the Randall-Sundrum scenario. (For related work see, e.g., [27].) Recall that in order to obtain  $Z_2$  symmetric domain-wall solutions of the

Randall-Sundrum scenario, the gauged supergravity potential would have to have two isolated supersymmetric *minima*. Since the potentials for the massless scalar fields in a gauged supergravity generically do not have this feature, we now turn to an alternative proposal [9] to include other scalar fields that do not lie in the massless supermultiplet.

We shall focus on the special classes of gauged supergravities that arise from sphere reductions of M-theory or string theory, with particular emphasis on the  $D = 5$  case. For examples in the Kaluza-Klein reduction of Type IIB string theory on a five-sphere ( $S^5$ ), there will be an infinite tower of massive supermultiplets in addition to the massless multiplet, and so one could consider the potentials for one or more of the massive scalar fields. In general, one cannot focus attention on a single such field in isolation, on account of its couplings to other fields. However, in certain special cases a consistent truncation to a single massive scalar can be performed. One such example is the “breathing mode” that parameterises the overall volume of the compactifying  $S^5$ . (Unlike the breathing mode in a toroidal reduction, which is massless, the breathing mode in a spherical reduction is a member of a massive supermultiplet.)

The scalar potentials for the breathing-mode scalars in various Kaluza-Klein spherical reductions were studied in [28]. Although the breathing mode is a member of a massive multiplet, the truncation is nonetheless consistent since it is a singlet under the isometry group of the internal sphere. (It would not in general be consistent to turn on a finite subset of other fields as well.)

The resulting  $D$ -dimensional Lagrangians all turn out to have the following form:

$$\mathcal{L}_D = R - \frac{1}{2}(\partial\phi)^2 - V, \quad (16)$$

where the potential is given by [28]

$$V = \frac{1}{2}g^2\left(\frac{1}{a_1^2}e^{a_1\Phi} - \frac{1}{a_1a_2}e^{a_2\Phi}\right). \quad (17)$$

The positive constants  $a_1$  and  $a_2$  are given by

$$a_1^2 = \frac{4}{N} + \frac{2(D-1)}{D-2}, \quad a_1 a_2 = \frac{2(D-1)}{D-2}, \quad (18)$$

where Type IIB reduction on  $S^5$  corresponds to  $N = 1$ ,  $D = 5$ . Since  $a_1 > a_2 > 0$ , the potential has a minimum at  $\phi = 0$ , corresponding to the self-dual point where the volume of the five-sphere and the radius of the  $\text{AdS}_5$  are equal (in appropriate units). In addition, this potential can be cast in the standard supersymmetric form:

$$V = \left(\frac{\partial W}{\partial\phi}\right)^2 - \frac{D-1}{2(D-2)}W^2, \quad (19)$$

where

$$W = \sqrt{\frac{N}{2}} g \left( \frac{1}{a_1} e^{a_1 \Phi/2} - \frac{1}{a_2} e^{a_2 \Phi/2} \right). \quad (20)$$

Thus, there is a domain wall solution that can be obtained in terms of the coupled first-order differential equations (15). Solving for  $\phi$  and  $A$  one finds the result:

$$A^{\frac{D-1}{2}} = c \frac{\partial W}{\partial \Phi} e^{-\frac{1}{2}(a_1+a_2)\Phi}, \quad (21)$$

where  $c$  is an integration constant and the solution for  $\Phi$  is given by:

$$z - z_0 = \frac{4}{a_2 g \sqrt{N}} e^{-\frac{1}{2}a_2 \Phi} {}_2F_1\left[\frac{a_2}{a_2 - a_1}, 1, 1 + \frac{a_2}{a_2 - a_1}; e^{\frac{1}{2}(a_1 - a_2)\Phi}\right]. \quad (22)$$

The solutions above *have two different branches*. In one branch,  $\phi$  runs from 0 to  $+\infty$ , with  $z$  running from  $z = -\infty$  to  $z = 0$ , where we have chosen the integration constant  $z_0$  to yield the following result:

$$\begin{aligned} e^{-\frac{1}{2}a_1 \phi} &\sim -\frac{1}{4}a_1 \sqrt{N} g z, \\ A^{\frac{(D-1)}{2}} &\sim c \sqrt{\frac{N}{8}} g e^{-\frac{1}{2}a_2 \phi} \sim c \sqrt{\frac{N}{8}} g \left( -\frac{1}{4}a_1 \sqrt{N} g z \right)^{\frac{a_2}{a_1}}. \end{aligned} \quad (23)$$

In this branch, when the coordinate  $z$  reaches its limit at  $z = 0$ , the metric factor therefore goes to zero, and there is a power-law naked curvature singularity. (Note that in this regime the solution extends into large positive values of the potential (17) with a large cost to the energy density of the wall, and it thus terminates at a finite value of the transverse coordinate.) As  $z$  approaches  $-\infty$ , the metric asymptotically approaches the AdS spacetime, described in horospherical coordinates with  $z \rightarrow -\infty$  corresponding to the Cauchy horizon. Note that on that side of the wall the gravity is repulsive and provides “one half” of the Randall-Sundrum wall.

The study of the gravitational fluctuating modes, internal to the wall, reduces to the study of the Schrödinger equation, whose potential [10] has an attractive, singular region near the naked singularity. Nevertheless the spectrum has energy levels bounded from below. However, the boundary conditions at the naked singularity exclude the massless normalizable mode (corresponding to the four-dimensional graviton). On the other hand corrections (of the order of the inverse string scale), that would smooth out the naked singularity, would in turn provide the non-singular attractive Schrödinger potential with precisely one normalizable massless state. Further investigations to identify the origin of the smoothing out of such singularities within the string theory context is needed.

In the second branch,  $\phi$  runs from 0 to  $-\infty$ , while  $z$  runs from  $z = -\infty$  to  $z = +\infty$ . The behaviour of the solution near  $z = -\infty$  is the same as in the branch discussed



previously, with the metric approaching asymptotically AdS. As  $z$  approaches  $+\infty$ , the solution becomes

$$\begin{aligned} e^{-\frac{1}{2}a_2\phi} &\sim \frac{1}{4}a_2\sqrt{N}gz, \\ A^{\frac{(D-1)}{2}} &\sim -c\sqrt{\frac{N}{8}}ge^{-\frac{1}{2}a_1\phi} \sim -c\sqrt{\frac{N}{8}}g\left(\frac{1}{4}a_2\sqrt{N}gz\right)^{\frac{a_1}{a_2}}. \end{aligned} \quad (24)$$

(The constant  $c$  is negative in this case.) This side describes one-side of a supersymmetric *dilatonic* domain wall. Interestingly, it has no curvature singularity; as  $z$  tends to  $+\infty$  the curvature falls off as  $1/z^2$ , while the diverging dilaton  $\phi \rightarrow -\infty$  approaches the weak coupling limit. Unfortunately, the dilatonic vacuum side does not provide for the gravity trapping solution.

Thus within a pure field-theoretic framework, i.e. by employing only the breathing-mode scalar field to construct the domain wall solution, one did not fully succeed in constructing domain wall solutions that would allow for trapping of gravity on the wall, though the first branch would provide for such a scenario if the mechanism to smooth out the naked singularity within string theory existed. (For a possible related mechanism see [29].)

In [9] it was therefore proposed to add a delta-function source to the effective Lagrangian. In this case, the diverging behaviour of the dilaton is cut-off by a delta-function source at some finite value of  $z$ , say  $z = z_*$  and the solution for  $z > z_*$  can be replaced by a reflection of the solution for  $z < z_*$ . Now the solution is  $Z_2$  symmetric and the metric factor  $A(z)$  falls-off on either side of the infinitely thin wall, supported by the delta function source, thus reproducing the Randall-Sundrum scenario. The origin of such a delta function source within Type IIB supergravity was further explored in [30].

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